

Inflation as a probe of new physics

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Abstract

In this paper we consider inflation as a probe of new physics near the string or Planck scale. We discuss how new physics can be captured by the choice of vacuum, and how this leads to modifications of the primordial spectrum as well as the way in which the universe expands during inflation. Provided there is a large number of fields contributing to the vacuum energy – as typically is expected in string theory – we will argue that both types of effects can be present simultaneously and be of observational relevance. Our conclusion is that the ambiguity in choice of vacuum is an interesting new parameter in serious model building.

1 Introduction

Is there a simple and unique theory describing the origin of the universe? To some extent the answer is essentially yes. The key is inflation, which provides a model for the initial conditions of the universe largely insensitive to the details of physics near the Planck scale. Whatever happened at the very first instance, the universe would, according to inflation, end up roughly in the same state as the one in which we find it today. An excellent introduction to the subject of inflation with references can be found in [1]. Nevertheless, it is reasonable to look further and find the reasons behind inflation and why it took place. String theory has recently reached a sufficient state of maturity to be able to address these issues and an intriguing new scenario is developing. Far from starting off in a unique state, an enormous range of possibilities has been discovered. Some can be described by known building blocks like branes and fluxes, but one can be sure that there are still others to be revealed. Many authors take this as a sign that chance played an important role in selecting our universe. The conclusion – for some depressing and for others stimulating – seems to be that the universe is, and always has been, a messy place.

Most attempts to describe the early universe use classical effective theory where quantum effects never play a prominent role. Recent notable exceptions, where true quantum cosmology is considered, can be found in e.g.[2][3][4]. The reasons are obvious, these are the situations where reliable calculations can be performed. But there is very little reason to believe that we will be able to get all the way without the full force of string theory and quantum gravity. In particular, having the above mentioned change in view of the origin of the universe in mind where chance rather than simplicity is the guide, it is natural to expect that whatever new effects we can imagine are also likely to be present and of relevance if we look hard enough.

We will therefore consider a different line of approach where we temporarily drop the hope of performing well controlled calculations. Instead we will look for qualitatively new types of phenomena but limit ourselves to rough estimates of their magnitude. The objective is to find examples where the estimates suggest signatures within observational reach. In doing so we will be pushing the limits of how insensitive inflation really is to unknown physics at higher energies or earlier times. To do this we assume that all effects due to unknown high energy quantum gravity or stringy physics can be captured by a choice of vacuum. This assumption is extremely natural in an expanding universe as has been spelled out in [5]. To understand why, consider, for simplicity, a massless field with modes that redshift as the universe expands. The key to the argument is that any given field mode can not reliably be traced further back in time than to an era when the wavelength was of order the fundamental length scale. At this point, as argued in [5], we parametrize our ignorance of what took place at even earlier times and smaller length scales through the choice of vacuum. The picture we have in mind is one where the mode emerges out of the space time foam in a particular state determined by high energy physics. All our calculations, however, will take place at low energies with crucial initial conditions imposed not at a fixed

time but at a fixed scale.

As argued in [5] and [6] there is a natural expectation on the range of vacua that can be expected and this leads to definite predictions of the kind of effects one should look for. If we had been allowed to follow a given mode arbitrarily far back in time we could have argued for a unique vacuum, the Bunch-Davies vacuum, which also is the vacuum typically used in calculations of the primordial spectrum. In the presence of a fundamental scale there is an ambiguity characterized by the ratio of the fundamental scale to the Hubble scale. It is this ambiguity that new physics beyond the fundamental scale can exploit to influence physics at lower scales.

The first example on how this could happen is concerned with the possibility of remaining traces of elusive high energy physics in the CMBR fluctuations or in large scale structures. The basic idea is that the inflaton, whose quantum fluctuations are responsible for the fluctuations in the CMBR, is probing physics near the string or Planck scale through the choice of vacuum as explained above. Arguments have been put forward that the natural size of these effects are just on the border of what might be detectable by planned CMBR-measurements, [7]. A positive detection would open a direct window to string theory and quantum gravity. The literature on the subject is rich – the list of references, [5-27], only include a selected few of all the papers written on the subject.

A non-trivial vacuum choice opens up the issue of back reaction on the geometry, which will be our second example on how new physics could lead to observational signatures. Typically, the issue of quantum back reaction is neglected and assumed to be part of the unsolved problem of the cosmological constant. In the cases we will be considering the nontrivial dependence of the vacuum energy on the cosmological parameters make the issue less clear cut. In fact, the problem was early realized in [28], where it was correctly concluded that the effects in general were too small to cause any problem. In [26] and [27] it was argued that far from being a problem, the presence of the vacuum contribution is instead an intriguing possibility. The vacuum energy will not, it was argued, in anyway spoil inflation but can under the right conditions contribute in a constructive way to the creation of an inflationary phase. There is even, as argued in [27], an automatic slow roll built into the model, which could be of observational interest.

In this paper we will investigate the possibility of a model that both has detectable modulations in the CMBR-spectrum and quantum vacuum energy which interferes with inflation. We begin by a short review of the essential results for the amplitude of the modulation in section 2, and in section 3 we review the issue of back reaction. In section 4 we put the pieces together and we end by some conclusions and outlook.

2 A modulated spectrum

We begin by considering the effect on the CMBR, or the primordial spectrum in general, due to new physics modelled by a choice of vacuum different from the usual

Bunch-Davies vacuum. Interestingly, the effects are quite generic. According to the analysis of [5], the typical effect to be expected on the primordial spectrum is of the form

$$P(k) = \left(\frac{H}{\dot{\phi}}\right)^2 \left(\frac{H}{2\pi}\right)^2 \left(1 - \frac{H}{\Lambda} \sin\left(\frac{2\Lambda}{H}\right)\right), \quad (1)$$

where we note a characteristic, relative amplitude of the correction given by $\frac{H}{\Lambda}$, and a modulation sensitively dependent on how $\frac{\Lambda}{H}$ changes with k . Λ is the energy scale of the new physics which could be the string scale or the Planck scale. The overall factor in front of the expression is the well known amplitude for the primordial spectrum due to fluctuations in the inflaton field. It is easy to qualitatively understand the presence of the modulation. Just as there are oscillations taking place after the fluctuations re-enter through the horizon, giving rise to the acoustic peaks, there are oscillations taking place before the fluctuations exit and freeze. These fluctuations are only present if we make a non-trivial choice of state for the inflaton different from the Bunch-Davies vacuum. The claim is that whatever the nature of the high energy physics really is, a modulated spectrum of the above described form is to be expected.

In [7] one can find a discussion of the phenomenological relevance of this effect and how the magnitude is related to the characteristic parameters describing the inflationary phase. Using the standard slow roll approximation, assuming a single field inflaton with potential $V(\phi)$, one has that the amplitude of the primordial spectrum is given by

$$\left(\frac{H}{\dot{\phi}}\right)^2 \left(\frac{H}{2\pi}\right)^2 = \frac{1}{24\pi^2 M_{pl}^4} \frac{V}{\varepsilon}, \quad (2)$$

where

$$\varepsilon = \frac{M_{pl}^2}{2} \left(\frac{V'}{V}\right)^2, \quad (3)$$

is one of the slow roll parameters. According to measurements of the CMBR we have

$$\frac{V^{1/4}}{\varepsilon^{1/4}} \sim 0.027 M_{pl}, \quad (4)$$

implying a relation between the Hubble constant and the slow roll parameter according to

$$\frac{H}{M_{pl}} \sim 4 \cdot 10^{-4} \sqrt{\varepsilon}. \quad (5)$$

We now assume initial conditions imposed at $\Lambda = \gamma M_{pl}$, which implies

$$\frac{\Delta k}{k} \sim \frac{\pi H}{\varepsilon \gamma M_{pl}} \sim \frac{\pi \beta^2}{\sqrt{3} \gamma \sqrt{\varepsilon}} \sim 1.3 \cdot 10^{-3} \frac{1}{\gamma \sqrt{\varepsilon}}, \quad (6)$$

and

$$\frac{H}{\Lambda} \sim 4 \cdot 10^{-4} \frac{\sqrt{\varepsilon}}{\gamma}. \quad (7)$$

These latter two relations are what we need in order to produce our estimates. In order to beat cosmic variance and get an effect which has a chance to be detected we need $\frac{H}{\Lambda} \sim 10^{-2}$, or, in other words, $\frac{\sqrt{\varepsilon}}{\gamma} \sim 20$. In order for the modulation to be detectable in the CMBR we would furthermore like to have $\frac{\Delta k}{k} \sim \mathcal{O}(1)$ and therefore $\varepsilon \sim 10^{-2}$. Remarkably, this is perfectly consistent with what one would expect from a generic heterotic string compactification with high string scale. The conclusion in [7], that the effects may be detectable by the upcoming Planck satellite, has essentially been verified in subsequent work such as [25]. The latter represents the most complete analysis to date.

But what if we are interested in a considerably lower string and Hubble scale? This is highly relevant for many of the brane universe scenarios. We might still arrange for a suitable value of the amplitude – if both H and the string scale are small – but with the accompanying small value of ε , $\frac{\Delta k}{k}$ will in general be much too large for observable modulations. Since the range of observationally relevant e-foldings, when large scale structure is included, is about 10, we can accept $\frac{\Delta k}{k} \sim \mathcal{O}(10)$ but not much more. In this case we might push down to $\varepsilon \sim 10^{-3}$, but this seems to be the limit.

On the other hand, with a sufficiently large $\frac{\Delta k}{k}$ we can raise the amplitude much further without getting in conflict with the present measurements of the CMBR. For instance, we can stick to $\varepsilon \sim 10^{-2}$ and consider an amplitude at the 10% level keeping $\frac{\Delta k}{k} \sim \mathcal{O}(10)$. This is an intriguing possibility where we expect a non-trivial mixing between effects due to the spectral parameter and effects due to the modulation. In case of an observationally detected running of the spectral parameter, when comparing the CMBR results with measurements of large scale structure, a modulation of the form discussed here could be of relevance.

This is not the whole story, however, As we will see below there is an interesting way to relax the constraints. We will return to this issue after investigating how the non-trivial vacuum might backreact in the expansion.

3 A modified expansion

As explained in the introduction, the presence of a nontrivial vacuum, motivated by the presence of unknown high energy physics, raises the issue of backreaction. Clearly, this is related to the problem of the cosmological constant and it is not obvious how this decouples from the issue of inflation itself. How can a subtraction of quantum fluctuations be argued for? Focusing on the contribution to the vacuum energy coming from the non-standard vacuum, as compared with the Bunch-Davies vacuum, we find an additional energy density given by

$$\rho_{\Lambda} \sim \frac{1}{2\pi^2} \int_0^{\Lambda} dp p^3 \frac{H^2}{\Lambda^2} = \frac{\Lambda^2 H^2}{8\pi^2}. \quad (8)$$

Clearly, this contribution to the vacuum energy can be neglected, to lowest order, as long as $\Lambda \ll M_p$. This was the conclusion of [28]. However, even if the effect is small it can affect the slow roll parameters and the way the expansion of the universe changes. In order to see this we must be a little bit more careful in our analysis.

Since H will be changing with time, i.e. decrease, we must take this into account when calculating the vacuum energy density. Modes with low momenta were created at earlier times when the value of H were larger, and there will be an enhancement in the way these modes contribute to the energy density. We therefore write

$$\rho_\Lambda(a) = \frac{1}{2\pi^2} \int_\varepsilon^\Lambda dp p^3 \frac{H^2\left(\frac{ap}{\Lambda}\right)}{\Lambda^2} = \frac{1}{2\pi^2} \frac{\Lambda^2}{a^4} \int_{a_i}^a dx x^3 H^2(x), \quad (9)$$

where we have introduced a low energy cutoff corresponding to the energy at the time of observation of modes that started out at Λ at some arbitrary initial scale factor a_i .¹

If we take a derivative of the energy density with respect to the scale factor and use $\frac{d}{da} = \frac{1}{aH} \frac{d}{dt}$, we find

$$\dot{\rho}_\Lambda + 4H\rho_\Lambda = \frac{1}{2\pi^2} \Lambda^2 H^3, \quad (10)$$

and we conclude that we must include a source term. With additional matter present, with an equation of state of the form

$$\dot{\rho}_m + 3H(\rho_m + p_m) = 0,$$

it was found in [27] that the evolution is governed by

$$\frac{d}{da} (a^5 H H') = -\frac{8\Lambda^2}{3\pi M_p^2} a^3 H^2 - \frac{4\pi}{M_p^2} a^3 (1 + w_m) (1 - 3w_m) \rho_m, \quad (11)$$

where we let $' = \frac{d}{da}$. The above way of expressing the content of the Friedmann equations might seem unfamiliar, but is really the natural way in the presence of the source term.

There are two constants of integration in this framework, as indicated in the solution given by

$$H^2 = C_1^2 a^{-2n_1} + C_2^2 a^{-2n_2} + \frac{8\pi}{3M_p^2} \frac{(1 + w_m)(1 - 3w_m)}{(1 + w_m)(1 - 3w_m) - \frac{16\Lambda^2}{9\pi M_p^2}} \rho_m, \quad (12)$$

where

$$n_{1,2} = 1 \pm \sqrt{1 - \frac{4\Lambda^2}{3\pi M_p^2}}. \quad (13)$$

The first of the constants of integration corresponds to how much radiation there is, the second one is an analogue of the cosmological constant. In accordance with this

¹Note that we are using another convention for Λ as compared to the paper [27] in order to facilitate the comparison with the expressions for the modulation.

we note the the last term in (11) and (12) vanishes in case of matter with an equation of state like the one of radiation or a cosmological constant. The reason is that these types of matter are already taken care of by the constants of integration, and become part of the initial conditions. There is an interesting and important difference from the sourceless case, though: the cosmological constant is rolling. As argued in [27] this could have interesting implications for the problem of the cosmological constant, but we will not discuss this issue further in this paper.

In [27] it was argued that the modification of the cosmological evolution due to the source did not spoil inflation. In fact, without the matter contribution, it gives rise to a natural slow roll with a slow roll parameter easily read off from (11) as

$$\varepsilon = \frac{2\gamma^2}{3\pi}, \quad (14)$$

for small $\gamma = \frac{\Lambda}{M_p}$.

The analysis has so far dealt with a particularly natural example of how the source term depends on the expansion of the universe. For completeness we can generalize the above analysis to a more general dependence on the Hubble constant. Assuming

$$\rho_\Lambda(a) = \frac{1}{2\pi^2} \int_\varepsilon^\Lambda dp p^3 g\left(\frac{H\left(\frac{ap}{\Lambda}\right)}{\Lambda}\right) = \frac{1}{2\pi^2} \frac{\Lambda^4}{a^4} \int_{a_i}^a dx x^3 g\left(\frac{H(x)}{\Lambda}\right) \quad (15)$$

where $g = h_n f^n$, $f = f(z) = \frac{H^2}{\Lambda^2}$, and $z = a^{-4}$, we find

$$f'' = -\frac{\Lambda^2}{3\pi M_p^2} h_n \frac{1}{z^2} f^n. \quad (16)$$

This is a differential equation of the form of Emden-Fowler. Apart from being exactly solvable for $n = 1$, as used above, one can also write down the solution for $n = 0$:

$$H^2 = C_1 a^{-4} + C_2 - \frac{4\Lambda^2}{3\pi M_p^2} h_n \ln a. \quad (17)$$

Again we find a an effect which can be interpreted as a rolling cosmological constant.

The question that is our main concern in this paper is whether we can find a non-trivial interplay between the vacuum energy driven expansion and the usual inflaton. In particular we will see whether we can relax the constraints considered in the previous section and find interesting effects also in brane universe models.

4 Putting things together

Is there a way to combine the two effects in the previous sections so as to make them both observationally interesting? We saw above that it was possible to achieve a

slow roll inflationary phase by using the backreaction from a non-standard vacuum. A problem with the present simple model is, however, that there is no natural way for inflation to end. For this we consider a model where in addition to the quantum vacuum energy there also is an inflaton. The inflaton takes the role of the matter field of the previous section, with a slow roll version of its equation of motion given by

$$3aH^2\phi' = -\frac{dV}{d\phi}. \quad (18)$$

The evolution equation for the Hubble constant now takes the form

$$\frac{d}{da} (a^5 H H') = -\frac{8n\Lambda^2}{3\pi M_p^2} a^3 H^2 - \frac{4\pi}{M_p^2} \frac{d}{da} \left(a^4 (aH\phi')^2 \right), \quad (19)$$

where we have generalized to an arbitrary number of fluctuating fields, n , all in the non-standard vacuum. As pointed out in [13], it is reasonable to expect that all fields, not only the inflaton, is being influenced by the unknown high energy physics in a similar way. We will, for simplicity, assume that all of these fields are effectively massless.

There will in general be a complicated interplay between the two terms, but we will focus on situations where one of them dominates. It will also be important to figure out when one term takes over after the other. Let us consider a couple of explicit examples. We start with a potential of the type of chaotic inflation given by

$$V = \frac{1}{2}m^2\phi^2. \quad (20)$$

To this we should add the contribution from the vacuum energy, but since this show up as a rolling cosmological constant through constants of integration, we do not put it in explicitly here. We can easily solve (18) with the result

$$\phi = \phi_0 a^{-\frac{m^2}{3H^2}}, \quad (21)$$

and the evolution equation becomes

$$\frac{d}{da} (a^5 H H') = -\frac{8n\Lambda^2}{3\pi M_p^2} a^3 H^2 - \frac{16\pi}{M_p^2} \phi^2 \frac{m^4}{H^4} a^3 H^2. \quad (22)$$

The first term alone would suggest a slow roll parameter given by

$$\varepsilon = \frac{2n\gamma^2}{3\pi},$$

while the second on its own would give

$$\varepsilon_{\text{inf}} = \frac{4\pi}{9} \frac{\phi^2}{M_p^2} \frac{m^4}{H^4}. \quad (23)$$

We will be interested in a situation where the roll is dominated by the vacuum energy, that is $\varepsilon \gg \varepsilon_{\text{inf}}$, during the era when the relevant fluctuations are leaving the horizon. As time proceeds, H will eventually come down far enough for ε_{inf} to dominate and eventually put an end to inflation. Let us see what kind of restrictions we can derive for the various parameters.

When it comes to the amplitude of the primordial spectrum the inflaton plays a crucial role. It is the inflaton that determines when inflation is going to end. The fluctuations calculated on a spatially flat surface are given by $\delta\phi$, but on a comoving surface, determined by a fixed value of ϕ , one finds curvature perturbations given by $\mathcal{R} = \frac{H\delta\phi}{\dot{\phi}}$, which remain constant after horizon exit. One might possibly worry that this is no longer true since we have played around with the conservation laws by introducing a source term. As argued in [27], however, the effect of the vacuum energy can be mimicked by matter with a peculiar equation of state obeying the usual conservation laws. In this way, we see that the overall amplitude as well as the amplitude of the modulations are governed by ε_{inf} – even though $\varepsilon_{\text{inf}} \ll \varepsilon$. Hence we have

$$\frac{H}{\Lambda} \sim 4 \cdot 10^{-4} \frac{\sqrt{\varepsilon_{\text{inf}}}}{\gamma}, \quad (24)$$

and in case of the wave length of the modulation we find

$$\frac{\Delta k}{k} \sim \frac{\pi H}{\varepsilon \gamma M_{\text{pl}}} \sim 1.3 \cdot 10^{-3} \frac{\sqrt{\varepsilon_{\text{inf}}}}{\gamma \varepsilon}, \quad (25)$$

where also ε appears. Interestingly, we now have a decoupling between the constraints thanks to the different roles played by ε_{inf} and ε . Arguing like in section 2 we find that these requirements are satisfied if $\frac{\sqrt{\varepsilon_{\text{inf}}}}{\gamma} \sim 20$ (that is $\varepsilon_{\text{inf}} \sim 400\gamma^2$) and $\varepsilon \sim 10^{-2}$. We see that in order to have the effect of the quantum vacuum energy to be comparable to the effect of the inflaton, we need a large number of fields, $n \sim 10^3$. If we want the quantum vacuum energy to dominate, we need n to be even larger. Such numbers are nevertheless not unreasonable in GUT or string theories. We furthermore note that the larger n is, the smaller Hubble constant and string scale is allowed.

After the era when the observationally relevant perturbations leave one can typically allow for only up to 60 e-foldings. The total number of e-foldings, including those taking place earlier, can certainly be much larger. It is crucial to check whether the inflaton is able to stop inflation fast enough. In other words, will ε_{inf} increase fast enough? Expressing the dependence on the scale factor in (23) through $a = e^N$, where N is the number of e-foldings and putting $a = 1$ when the fluctuations leave the horizon, we find

$$\varepsilon_{\text{inf}} \sim \frac{4\pi}{9} \frac{\phi_0^2}{M_p^2} \frac{m^4}{H_0^4} e^{4N\varepsilon - \frac{2m^2}{3H_0^2} N e^{2N\varepsilon}}. \quad (26)$$

We see that ε_{inf} starts off increasing with N , reaches a maximum and then decreases again. The crucial issue is whether the maximum value is large enough and can be reached fast enough. Unfortunately, with values as small as $N = 60$ and $\varepsilon = 10^{-2}$ we can not achieve an ε_{inf} which starts off well below ε and then manages to take

over simply because $e^{4N\varepsilon - \frac{2m^2}{3H_0^2}Ne^{2N\varepsilon}} < e^{4N\varepsilon} \lesssim 10$. There is, nevertheless, room for an interesting interplay with contributions of the same order.

We next turn to our second example where we assume a potential of the opposite sign,

$$V = -\frac{1}{2}m^2\phi^2, \quad (27)$$

where again we should remember that there also is a contribution from the vacuum energy. The expression for ε_{inf} is the same in terms of ϕ , but we now have

$$\phi = \phi_0 a^{\frac{m^2}{3H^2}}, \quad (28)$$

with the inflaton rolling in the opposite direction, and

$$\varepsilon_{\text{inf}} \sim \frac{4\pi}{9} \frac{\phi_0^2}{M_p^2} \frac{m^4}{H_0^4} e^{4N\varepsilon + \frac{2m^2}{3H_0^2}Ne^{2N\varepsilon}}. \quad (29)$$

The situation is now slightly better since we are guaranteed an end to inflation if we just wait long enough. The numerics work out such that with $\frac{m^2}{H_0^2} \sim 0.02$ and ϕ_0 a little less than the Planck scale, inflation ends within the required 60 e-foldings.

Similar results can be obtained using other potentials, but these examples suffice to show that there is room for interesting interplay between the two sources of rolling, provided we have a large number of degrees of freedom contributing to the vacuum energy. An important observation is that the consistency relation for the size of the tensor contributions is violated in these models. That is, even if the roll is not that slow, we still can have the tensor perturbations heavily suppressed. The reason is the decoupling between ε_{inf} which governs the roll of ϕ , and ε which governs the roll of H (if $\varepsilon_{\text{inf}} \ll \varepsilon$). This is a property common to multifield inflationary models where the contribution from tensor modes also is suppressed. Interestingly, things work out in a similar way for holographic inflation, [31] (for a recent paper with references see [32]). In the case of holographic inflation there is also a modification of the Friedmann equations leading to new behavior, and just as in our model one needs a large number of fields to get an interesting effect, [33].

5 Conclusions and speculations

We have seen how new physics captured by a nontrivial vacuum choice can lead to generic effects on the primordial spectrum and the expansion of the early universe. There is even a possibility of having both effects present provided there is a large number of fields contributing to the quantum vacuum energy. But how can we embed this phenomenological approach within string theory? How do we make real calculations?

In a realistic scenario we would expect corrections to the potential and it is reasonable to let those determine the physics of the end of inflation. In fact, the models

used for brane inflation can more or less be taken over directly. In models such as those discussed in [29] and [30] the potential takes the form (27) for small values of ϕ , but then rapidly gets corrections as the branes move closer. In this way it is easy to get an ε_{inf} which dominates within the required 60 e-foldings. It is far from clear, however, how to obtain explicit constraints on the various parameters. How is the non-trivial vacua to be understood from the higher dimensional string theory point of view? In the brane universe scenarios the main contribution to the vacuum energy driving inflation comes from the tension of the branes, adjusted by warp factors. How is this changed in the presence of the rolling due to vacuum fluctuations? The difficulties are a reflection on our inability to address the problems of quantum gravity in a realistic setting. The answers should, clearly, be given by string theory and it would be interesting to incorporate the physics discussed in the present paper into a full-fledged brane universe model.

The effects we have studied influence the form of the primordial spectrum and interfere with the way the universe expands. We have argued that the magnitude of the effects are such that they need to be taken into account in inflationary model building. This suggests new possibilities for string theory and quantum gravity to become observationally relevant.

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